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Semileptonic Decays Of B Mesons Into Excited Charm Mesons: Leading Order And $1/m_c$ Contributions

Thomas Mannel

Institut für Kernphysik, Technische Hochschule Darmstadt

Schloßgartenstraße 9, 6100 Darmstadt, Germany

Winston Roberts

Department of Physics, Old Dominion University, Norfolk, VA 23529, USA

and

Continuous Electron Beam Accelerator Facility

12000 Jefferson Avenue, Newport News, VA 23606, USA.

We use the heavy quark effective theory to investigate the form factors that describe the semileptonic decays of a B meson into excited daughter mesons. For an excited daughter meson with charm, a single form factor is needed at leading order, while five form factors and two dimensionful constants are needed to order $1/m_c$ in the heavy quark expansion. For non-charmed final states, a total of four form factors are needed at leading order. For the process $B \to D^{(*)} X \ell \nu$, four form factors are also needed at leading order.

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I. INTRODUCTION

The heavy quark effective theory (HQET) [1] has been a very useful tool in the study of electroweak decays of heavy hadrons. It has been used to find relationships among the form factors describing the weak decays of heavy mesons, namely $B \to D$ and $B \to D^*$. It has also been applied to decays of heavy baryons [2] to leading order in the heavy quark Lagrangian, as well as to $1/m_Q$ correction terms in the decays of both heavy mesons and baryons [3].

In addition to these exclusive processes into the so-called elastic channels, Isgur and Wise [4] have studied decays into 'inelastic' channels, and have discussed their results in terms of a Bjorken sum rule [5], which has also been developed using the techniques of HQET. More recently, Falk [6] has developed representations of states of arbitrary spin.

In this paper, we consider another set of electroweak processes, namely the decay of the B meson into any of the allowed excited charm daughter mesons. The motivation here is that the B can decay weakly, but because the phase space available is sizeable, decays to excited daughter mesons (with the subsequent decays into multihadron final states) may provide a large contribution to the total decay rate of the B.

In fact, about twenty-three percent of charged B meson decays are to the inclusive channel $B \to \ell\nu +$ hadrons [7]. In addition, recent experiments indicate that the D and D^* channels, the so-called elastic channels, saturate the B semileptonic decays (to charm) to only about seventy percent [8]. This is evidence that semileptonic decays to excited mesons form a significant fraction of the total semileptonic decay rate of the B meson. Undoubtedly, decays to the lower lying excited mesons, such as those already treated by Isgur and Wise [4], will provide the major contribution to the inelastic channels. Nevertheless, in the event that semileptonic decays to even more highly excited states are observed, we enumerate the form factors necessary to describe such decays.

In addition to these decays, there is some interest in decays of the type $B \to D^{(*)}X\ell\nu$. In particular, if X is one or two pions, this costs relatively little in phase space, and may receive both resonant and non-resonant contributions. It may therefore be preferrable to treat these decays on a separate footing from the decays to charm resonances.

The corresponding decays of the D meson are, of course, more difficult to treat. Experimentally, the situation is probably more promising than in the decays of the B meson, since there is no expectation that the K and K^* channels should saturate the semileptonic decays of this meson. For this reason it may be even more interesting to examine semileptonic decays of D mesons to excited final states, despite the apparent theoretical difficulty.

In the next section of this note, we briefly discuss the representations of the states in which we are interested. We treat the semileptonic decays of the B meson to excited charm mesons to leading order, as well as the semileptonic decays of the D meson to excited mesons, and $B \to D^{(*)}X\ell\nu$, in section III. In section IV, we include $1/m_c$ terms in the semileptonic decays of the B meson to excited charm mesons, and in section V we present our conclusions.

II. REPRESENTATION OF $D^{(j,\ell)}$ STATES

An excited D meson with total angular momentum J will, in general, be represented by an object linear in a polarisation tensor, $\eta^{\mu_1\cdots\mu_J}(v)$. This polarization tensor is symmetric, transverse and traceless. The latter two properties are expressed by

$$v_{\mu_1}\eta^{\mu_1\cdots\mu_J}(v) = 0, \; g_{\mu_1\mu_2}\eta^{\mu_1\cdots\mu_J}(v) = 0. \tag{1}$$

For a state consisting of a heavy quark Q and a light component with the quantum numbers of an antiquark, the specific representation of any particular state will depend on the angular momentum j of the light component (antiquark) of the state. It is thus more convenient to refer to j than to J, since there will be a degenerate doublet of states with $J = j \pm 1/2$.

The parity of such a state is determined by the orbital angular momentum $\ell = j \pm 1/2$ between the light component and the heavy quark. For states with $\ell = j - 1/2$, Falk [6] writes the representations as

$$\begin{aligned}
V^{\mu_1\cdots\mu_k} &= \frac{1}{\sqrt{2}}\Lambda_+\eta^{\mu_1\cdots\mu_{k+1}}\gamma_{\mu_{k+1}}, \\
P^{\mu_1\cdots\mu_k} &= \sqrt{\frac{2k+1}{2k+2}}\Lambda_+\eta^{\nu_1\cdots\nu_k}\left[\delta^{\mu_1}_{\nu_1}\cdots\delta^{\mu_k}_{\nu_k} - \frac{1}{2j+1}A^{\nu_1\cdots\nu_k}_{\mu_1\cdots\mu_k}\right]\gamma_5, \\
A^{\nu_1\cdots\nu_k}_{\mu_1\cdots\mu_k} &= \left[\gamma_{\nu_1}\left(\gamma^{\mu_1}+v^{\mu_1}\right)\delta^{\mu_2}_{\nu_2}\cdots\delta^{\mu_k}_{\nu_k} - \cdots - \gamma_{\nu_1}\delta^{\mu_1}_{\nu_1}\cdots\delta^{\mu_{k-1}}_{\nu_{k-1}}\left(\gamma^{\mu_k}+v^{\mu_k}\right)\right], \\
\Lambda_+ &= \frac{1+\not{v}}{2}.
\end{aligned} \tag{2}$$

In the above, $k = j - 1/2$, and the state V has $J = j + 1/2 = \ell + 1$, while P has $J = j - 1/2 = \ell$. These two states are degenerate at leading order in HQET.

For light component with $\ell = j + 1/2$, the two possible states have the same values of J ($J = j \pm 1/2$) as the states described above. They are the counterparts of the V and P states, but with an extra factor of γ_5 to account for the different parity. This arises because ℓ for the (P, V) doublet of states differs by unity from ℓ for the (B, S) doublet of states with the same j. Thus, one may write

$$B^{\mu_1\cdots\mu_k} = \frac{1}{\sqrt{2}}\Lambda_+\eta^{\mu_1\cdots\mu_{k+1}}\gamma_{\mu_{k+1}}\gamma_5,$$

$$S^{\mu_1\cdots\mu_k} = \sqrt{\frac{2k+1}{2k+2}}\Lambda_+\eta^{\nu_1\cdots\nu_k}\left[\delta^{\mu_1}_{\nu_1}\dots\delta^{\mu_k}_{\nu_k} - \frac{1}{2j+1}A^{\nu_1\cdots\nu_k}_{\mu_1\cdots\mu_k}\right], \tag{3}$$

with $A^{\nu_1\cdots\nu_k}_{\mu_1\dots\mu_k}$ as defined above. We emphasize that, to leading order in HQET, the (P, V) states form a degenerate doublet, as do the (B, S) states, but the two doublets will, in general, have different masses. Note that for all four of these states, denoted $X^{\mu_1\cdots\mu_k}$,

$$\begin{aligned} v_{\mu_1}X^{\mu_1\cdots\mu_k} &= 0, \\ \gamma_{\mu_1}X^{\mu_1\cdots\mu_k} &= X^{\mu_1\cdots\mu_k}\gamma_{\mu_1} = 0, \\ \not{v}X^{\mu_1\cdots\mu_k} &= X^{\mu_1\cdots\mu_k}, \\ X^{\mu_1\cdots\mu_k}\not{v} &= \pm X^{\mu_1\cdots\mu_k}. \end{aligned} \tag{4}$$

In the last of these equations, the negative sign occurs for the (P, V) doublets of states, while the positive sign occurs for the (B, S) doublets of states. Further details of the structure and properties of these states are given in Falk's article [6].

III. $B \to D^{(j,\ell)}$: LEADING ORDER FORM FACTORS

To leading order in HQET, we are interested in the matrix element

$$\mathcal{A} \equiv < D^{(j,\ell)}(v')|\bar{h}^{(c)}_{v'}\Gamma h^{(b)}_v|B(v) > . \tag{5}$$

Representing the B meson by

$$B(v) \to \frac{1}{\sqrt{2}}\Lambda_+\gamma_5 \equiv M_B(v), \tag{6}$$

representing any of the states discussed in the previous section by $M^{\mu_1\cdots\mu_k}_{D^{(j,\ell)}}(v')$, and using HQET, we may write this matrix element as

$$< D^{(j,\ell)}(v')|\bar{h}^{(c)}_{v'}\Gamma h^{(b)}_v|B(v) >= \mathrm{Tr}\left[R_{\mu_1\dots\mu_k}\bar{M}^{\mu_1\cdots\mu_k}_{D^{(j,\ell)}}(v')\Gamma M_B(v)\right]. \tag{7}$$

$R_{\mu_1\dots\mu_k}$ is the most general tensor that we can build. Clearly, $M_{\mu_1\dots\mu_k}$ may contain no factors of v'_{μ_i}, nor $g_{\mu_i\mu_j}$, nor γ_{μ_i}. One is left with

$$\begin{aligned} R_{\mu_1\dots\mu_k} = v_{\mu_1}\dots v_{\mu_k}\Big[&\xi^{(j,\ell)}(v\cdot v') + \xi^{(j,\ell)\prime}(v\cdot v')\not{v} \\ &+\xi^{(j,\ell)\prime\prime}(v\cdot v')\not{v}' + \xi^{(j,\ell)\prime\prime\prime}(v\cdot v')\not{v}\not{v}'\Big]. \end{aligned} \tag{8}$$

Inspection shows that the $\not{v}$, $\not{v}'$ and $\not{v}\not{v}'$ terms are redundant, so that we may write, to this order,

$$\begin{aligned} < D^{(j,\ell)}(v')|\bar{h}^{(c)}_{v'}\Gamma h^{(b)}_v|B(v) > = v_{\mu_1}\dots v_{\mu_k}\xi^{(j,\ell)}_0(v\cdot v') \\ \times\mathrm{Tr}\left[\bar{M}^{\mu_1\cdots\mu_k}_{D^{(j,\ell)}}(v')\Gamma M_B(v)\right]. \end{aligned} \tag{9}$$

Thus, a single form factor is needed to this order, regardless of the spin of the final meson.

States with $J^P = 0^-, 1^-$, $j = 1/2$, $\ell = 0$ may be thought of as radial excitations of the ground states $D^{(*)}$. Because of the heavy quark symmetry, and the orthogonality of these states with respect to the ground states which have the same quantum numbers, we must have

$$\begin{aligned} < D^{(n,\ell=0)}(v')|\bar{h}^{(c)}_{v'}\Gamma h^{(b)}_v|B(v) > = (v\cdot v' - 1)\xi^{(n,1/2,0)}(v\cdot v') \\ \times\mathrm{Tr}\left[\bar{M}_{D^{(*)}}(v')\Gamma M_B(v)\right], \end{aligned} \tag{10}$$

where the superscripts (n) denote the nth radial excitation. That is, these amplitudes must vanish as $v' \to v$. Note, too, that all of the other amplitudes ($j \neq 1/2$) vanish trivially at the non-recoil point.

We end this section by noting that we can use similar methods to enumerate the form factors that occur in the corresponding heavy to light transitions. For such a transition, say $D \to K^{(n)}$, where $K^{(n)}$ is some excited light meson, we write

$$< K^{(n)}(p)|\bar{s}\Gamma h^{(c)}_v|D(v) >= \mathrm{Tr}\left[R(\eta, v, p)\Gamma M_D(v)\right], \tag{11}$$

where R is a Dirac matrix that must be linear in the symmetric, traceless, transverse, n-index, polarization tensor η that describes the daughter meson of spin n. On inspection, we obtain

$$\begin{aligned} R(\eta, v, p) \equiv v_{\mu_1}\dots v_{\mu_{n-1}}\eta^{*\mu_1\cdots\mu_n} \\ \times\left[v_{\mu_n}\left(\zeta^{(n)}_1 + \not{p}\zeta^{(n)}_2\right) + \gamma_{\mu_n}\left(\zeta^{(n)}_3 + \not{p}\zeta^{(n)}_4\right)\right]. \end{aligned} \tag{12}$$

Thus, four form factors are needed, in general, for each of these transitions. For the left handed current, $\Gamma = \gamma_\mu(1-\gamma_5)$, HQET has not simplified things much. For a scalar or pseudoscalar daughter meson, the ζ_3 and ζ_4 terms do not contribute, and two form factors are needed, as would be the case in the general treatment of such decays. For daughter hadrons of higher spin, four form factors are needed, but this is also the case in the general treatment of such decays. HQET only starts being advantageous, as far as the number of form factors go, when the matrix elements of currents other than the left handed current, such as $\Gamma = \sigma_{\mu\nu}$, for daughter mesons with spin greater than one, are considered, since for such combinations of currents and mesons, more than four form factors would be needed in the general treatment.

In the same manner, we may also write down the form of the matrix element needed for describing decays of the type $B \to D^{(*)} X \ell \bar{\nu}_\ell$, where X may be any hadronic state. We only require that the total spin n of the state X be known. This is akin to a partial wave analysis. Such a matrix element may be written as

$$< D^{(*)}(v') X(p) | \bar{h}_{v'}^{(c)} \Gamma h_v^{(b)} | B(v) > = \mathrm{Tr}\left[\Theta_X(\eta, p, v', v) \bar{M}_{D^{(*)}}(v') \Gamma M_B(v) \right], \tag{13}$$

where, again, Θ is a Dirac matrix that must also be linear in the symmetric, traceless, transverse, n-index, polarization tensor η, that describes the system X of spin n. Inspection yields

$$\begin{aligned} \Theta_X(\eta, p, v', v) = & \, w_{\mu_1} \ldots w_{\mu_{n-1}} \eta^{*\mu_1 \ldots \mu_n} \\ & \times \left[w_{\mu_n} \left(\aleph_1^X + \not{p} \aleph_2^X \right) + \gamma_{\mu_n} \left(\aleph_3^X + \not{p} \aleph_4^X \right) \right], \end{aligned} \tag{14}$$

where $w = v - v'$.

IV. $B \to D^{(j,\ell)}$: $\frac{1}{M_Q}$ TERMS

We write the $1/m_c$ part of the HQET Lagrangian as

$$\mathcal{L}' = \frac{1}{2m_c} \bar{h}_{v'}^{(c)} \left[D^\mu (g_{\mu\nu} - v'_\mu v'_\nu) D^\nu + \frac{g_s}{2} \sigma^{\mu\nu} F_{\mu\nu} \right] h_{v'}^{(c)}, \tag{15}$$

where the coefficients of the various terms are obtained from tree-level matching to the full QCD Lagrangian. In addition, tree-level matching of the local operator $\bar{h}_{v'}^{(c)} \Gamma h_v^{(b)}$ to order $1/m_c$ leads to

$$\bar{h}_{v'}^{(c)} \Gamma h_v^{(b)} \to \bar{h}_{v'}^{(c)} \left[\Gamma - \frac{i}{2m_c} \overleftarrow{\not{D}} \Gamma \right] h_v^{(b)}. \tag{16}$$

We begin with the local term by examining the matrix element

$$< D^{(j,\ell)}(v') | \bar{h}_{v'}^{(c)} \overleftarrow{D}_\lambda \hat{\Gamma} h_v^{(b)} | B(v) > = \mathrm{Tr}\left[S_{\mu_1 \ldots \mu_\ell \lambda} \bar{M}_{D^{(j,\ell)}}^{\mu_1 \ldots \mu_\ell}(v') \hat{\Gamma} M_B(v) \right]. \tag{17}$$

The most general form for $S_{\mu_1 \ldots \mu_\ell \lambda}$ is

$$\begin{aligned} S_{\mu_1 \ldots \mu_\ell \lambda} = & \, v_{\mu_1} \ldots v_{\mu_\ell} \left[\xi_+^{(j,\ell)} (v + v')_\lambda + \xi_-^{(j,\ell)} (v - v')_\lambda + \xi_3^{(j,\ell)} \gamma_\lambda \right] \\ & + \xi_4^{(j,\ell)} v_{\mu_1} \ldots v_{\mu_{\ell-1}} g_{\mu_\ell \lambda} \end{aligned} \tag{18}$$

If we multiply eqn. (17) by v'^λ, and use the equation of motion of $\bar{h}_{v'}^{(c)}$, we obtain

$$\xi_+^{(j,\ell)} (1 + v \cdot v') + \xi_-^{(j,\ell)} (v \cdot v' - 1) \pm \xi_3^{(j,\ell)} = 0. \tag{19}$$

Here, the positive sign on the $\xi_3^{(j,\ell)}$ term occurs for the $(B,\ S)$ doublets of states, while the negative sign occurs for the $(P,\ V)$ doublets. In addition, we can use conservation of momentum to write [3]

$$\begin{aligned} & < D^{(j,\ell)}(v') | i\partial_\lambda (\bar{h}_{v'}^{(c)} \Gamma h_v^{(b)}) | B(v) > \\ & = \left[(m_B - m_b) v_\lambda - (m_{D^{(j,\ell)}} - m_c) v'_\lambda \right] < D^{(j,\ell)}(v', f) | \bar{h}_{v'}^{(c)} \Gamma h_v^{(b)} | B(v) >. \end{aligned} \tag{20}$$

The matrix element on the right hand side is the leading order matrix element, described in terms of the single form factor $\xi_0^{(j,\ell)}$. Let us define $\Delta m_b \equiv m_B - m_b$, and $\Delta m_c \equiv m_{D^{(j,\ell)}} - m_c$. Multiplying eqns. (17) and (20) by v^λ, and using the equation of motion of $h_v^{(b)}$ leads to

$$\begin{aligned} (\Delta m_b - \Delta m_c v \cdot v') \xi_0^{(j,\ell)} = & \, \xi_+^{(j,\ell)} (1 + v \cdot v') + \xi_-^{(j,\ell)} (1 - v \cdot v') \\ & - \xi_3^{(j,\ell)} + \xi_4^{(j,\ell)}, \end{aligned} \tag{21}$$

which, when combined with eqn. (19), gives

$$2\xi_-^{(j,\ell)} (1 - v \cdot v') + \xi_4^{(j,\ell)} = (\Delta m_b - \Delta m_c v \cdot v') \xi_0^{(j,\ell)}, \tag{22}$$

for the $(P,\ V)$ doublets of states, and

$$2\xi_+^{(j,\ell)} (1 + v \cdot v') + \xi_4^{(j,\ell)} = (\Delta m_b - \Delta m_c v \cdot v') \xi_0^{(j,\ell)}, \tag{23}$$

for the $(B,\ S)$ doublets. When we consider $\hat{\Gamma} = \gamma^\lambda \Gamma$, the $\xi_4^{(j,\ell)}$ term of eqn. (18) vanishes.

From the $1/m_c$ part of the Lagrangian, we may ignore the $(v.D)^2$ term, since this vanishes because of the equation of motion of $\bar{h}_{v'}^{(c)}$. The D^2 term gives

$$
\begin{aligned}
&< D^{(j,\ell)}(v')|T\int d^4x\left(\bar{h}_{v'}^{(c)}D^2h_{v'}^{(c)}\right)(x)\left(\bar{h}_{v'}^{(c)}\Gamma h_v^{(b)}\right)(0)|B(v)> \\
&= \mathrm{Tr}\left[S'_{\mu_1\ldots\mu_\ell}\bar{M}_{D^{(j,\ell)}}^{\mu_1\ldots\mu_\ell}(v')\left(\frac{1+\not{v}'}{2}\right)\bar{\Gamma}M_B(v)\right] \\
&= v_{\mu_1}\ldots v_{\mu_\ell}\chi_1^{(j,\ell)}(v\cdot v')\mathrm{Tr}\left[\bar{M}_{D^{(j,\ell)}}^{\mu_1\ldots\mu_\ell}(v')\Gamma M_B(v)\right].
\end{aligned}
\tag{24}
$$

This renormalizes the leading order form factor.

Finally, we are left with the $\sigma^{\alpha\beta}$ term, which is

$$
\begin{aligned}
&\frac{g_s}{4m_c} < D^{(j,\ell)}(v')|T\int d^4x\left(\bar{h}_{v'}^{(c)}\sigma^{\alpha\beta}F_{\alpha\beta}h_{v'}^{(c)}\right)(x)\left(\bar{h}_{v'}^{(c)}\Gamma h_v^{(b)}\right)(0)|B(v)> \\
&= \mathrm{Tr}\left[T_{\mu_1\ldots\mu_\ell\alpha\beta}\bar{M}_{D^{(j,\ell)}}^{\mu_1\ldots\mu_\ell}(v')\sigma^{\alpha\beta}\left(\frac{1+\not{v}'}{2}\right)\bar{\Gamma}M_B(v)\right].
\end{aligned}
\tag{25}
$$

The most general form for $T_{\mu_1\ldots\mu_\ell\alpha\beta}$ is

$$
T_{\mu_1\ldots\mu_\ell\alpha\beta} = v_{\mu_1}\ldots v_{\mu_\ell}\left[\chi_2^{(j,\ell)}\sigma_{\alpha\beta} + \chi_3^{(j,\ell)}(\gamma_\alpha v_\beta - \gamma_\beta v_\alpha)\right]. \tag{26}
$$

To this order, then, we find that four new form factors are needed to describe the decay of the B into any $D^{(j,\ell)}$ resonance, together with two unknown dimensionful constants. However, one of these new form factors, χ_1, may be absorbed into the leading order form factor. We have expressed it as an independent form factor, since it is a term that explicitly breaks the heavy quark symmetry of HQET. These results are consistent with those obtained by Luke [3] for the $B \to D^{(*)}$ semileptonic decays.

V. CONCLUSION

We have applied the techniques of HQET to the semileptonic decays of B mesons to excited charmed mesons, both to leading order and to order $1/m_c$. We have also investigated the semileptonic decays $B \to D^{(*)}X\ell\nu$ at leading order. Both of these classes of processes are phenomenologically interesting, as they are expected to form a sizeable contribution to the total decay rate of the B meson. There is, in fact, some overlap between these two sets of processes, as the $D^{(*)}X$ system may well be produced from the strong decay of an excited charm meson.

In addition to the decays mediated by the $b \to c$ current, we have examined the decays $D \to K^{(n)}\ell\nu$, where $K^{(n)}$ is some excited strange meson. In fact, the $K^{(n)}$ denoted above need not be a resonance, as the representation of the light meson only takes account of the total spin of the final state, but not how that spin is distributed among the hadrons that comprise the final state. Thus, one

may use the form factors we have listed in section III to describe either resonant or non-resonant decays, with the proviso that the kinematic structure of the form factors will be dictated by the decay being described.

One possible application of this work would be to investigate the relationship between the form factors we have discussed herein, and the sum rules of Bjorken, Dunietz and Taron [9]. Of course, without a specific parametrization or model of these form factors, comparison with experiment is not possible, except at the non-recoil point, $v = v'$. Such a comparison is even more problematic in the decays of the D meson, since not even the 'non-recoil' point is of any theoretical significance, as the normalization of the form factors is not *a priori* known: one must rely on model calculations if one wants to say anything about such decays.

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